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A brief review of “little string theories”

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Abstract.

This is a brief review of the current state of knowledge on “little string theories”, which are non-gravitational theories having several string-like properties. We focus on the six dimensional maximally supersymmetric “little string theories” and describe their definition, some of their simple properties, the motivations for studying them, the DLCQ and holographic constructions of these theories and their behaviour at finite energy density. (Contribution to the proceedings of Strings '99 in Potsdam, Germany.)

1. Introduction

One of the most surprising results which came out of the developments in string theory in the last few years is the existence of consistent non-gravitational theories in five and six space-time dimensions, even though no consistent Lagrangians are known for interacting theories in these dimensions. Generally these theories have been discovered by considering some limit of string theory configurations involving 5-branes and/or singularities. The higher dimensional non-gravitational theories may be divided into two classes. One class, which generally arises from a low-energy limit of string theory (or M theory), includes superconformal field theories in five and six dimensions[‡]. These theories seem to be standard local field theories, even though they have no good Lagrangian description (they are sometimes called “tensionless string theories” since many of them have BPS-saturated strings on their moduli space whose tensions go to zero at the conformal point). We will not discuss these theories here. The other class of theories, which was given the name “little string theories” (LSTs) by [1], is generally obtained by taking the string coupling to zero in some configuration of NS 5-branes and/or singularities which is not well-described by perturbation theory, and which in fact remains non-trivial (but decoupled from gravity) even after taking the string coupling to zero. The string scale M_s remains constant in this limit and plays an important role in the dynamics of these theories; they appear to be non-local theories with some string theory-like properties. In this contribution I will try to summarize all that is currently

[‡] Supersymmetry plays an essential role in proving the existence of these theories, it is not clear if non-supersymmetric theories also exist above four dimensions or not.

known about the simplest theories of this type, which are six dimensional theories with 16 supercharges arising as the $g_s \rightarrow 0$ limit of NS 5-branes in type IIA or type IIB string theory. There are many interesting results on lower dimensional LSTs, LSTs with less supersymmetry and compactifications of LSTs, that I will not have room to review here. The transparencies and audio for this talk are available at [2].

1.1. Definition and simple properties of “little string theories”

The simplest definition of six dimensional LSTs with 16 supercharges comes from looking at k parallel and overlapping NS 5-branes in type IIA or type IIB string theory, with $k > 1$, and taking the string coupling $g_s \rightarrow 0$ [3] (see also [4, 5]). The gravitational interactions go to zero in this limit, but the couplings of the fields associated with the NS 5-branes remain non-trivial even after taking g_s to zero; this can be seen, for instance, by analyzing the low-energy theory (as described below). The string scale M_s is kept finite in the limit, and it is the only parameter of the LSTs (except for the discrete parameter k); in particular, there is no continuous dimensionless coupling parameter in these theories, so unlike conventional string theories they have no obvious weakly coupled limits which can serve as the starting point for a perturbative expansion. The NS 5-branes break half of the supersymmetry of type II string theories, and thus there are no forces between them and we can indeed put k of them on top of each other. In six dimensions (like in ten or two dimensions) supersymmetry is generally chiral, and the dimensional reduction of type II string theory to six dimensions has $\mathcal{N} = (2, 2)$ supersymmetry. In the type IIA case the NS 5-branes preserve a chiral half of the supersymmetry so the resulting LST has $\mathcal{N} = (2, 0)$ supersymmetry, while in the type IIB case it has $\mathcal{N} = (1, 1)$ supersymmetry.

We can obtain equivalent definitions of the LSTs by using various duality symmetries. Using the duality between type IIA string theory and M theory we see that the $(2, 0)$ LST can also be derived from k 5-branes in M theory, with a transverse circle of radius R , in the limit $R \rightarrow 0$, $M_p \rightarrow \infty$ with $RM_p^3 = M_s^2$ kept constant. Using S-duality in type IIB we see that the $(1, 1)$ LST can be derived from the $g_s \rightarrow \infty$ limit of k D5-branes in type IIB string theory.

Additional definitions arise from recalling that k NS 5-branes with a transverse circle are T-dual to an A_{k-1} singularity with a compact circle [6], and noting that the existence of such a circle does not affect the decoupled theory on the NS 5-branes (it affects only the bulk modes which decouple when $g_s \rightarrow 0$). Thus, by using this T-duality, we conclude that the $(2, 0)$ LSTs arise also as the $g_s \rightarrow 0$ limit of type IIB string theory on an A_{k-1} singularity, and the $(1, 1)$ LSTs arise also as the $g_s \rightarrow 0$ limit of type IIA string theory on an A_{k-1} singularity[§]. The last definition suggests a generalization to the $g_s \rightarrow 0$ limit of D_n and E_n type singularities, which also correspond to LSTs with

[§] There is a subtlety here involving the degree of freedom corresponding to the center of mass position of the NS 5-branes which is not evident in the T-dual picture, generally we can ignore this degree of freedom since it is free and decoupled.

16 supercharges; we will not discuss these theories in detail here, but let us note that the full classification of LSTs in six dimensions with 16 supercharges includes theories with $(2,0)$ and $(1,1)$ supersymmetry of type G (where $G = A_k, D_k$ or E_k) for any simply-laced Lie algebra G .

Some simple properties of the LSTs with 16 supercharges follow from the above definitions :

- (i) Low-energy behaviour : For the $(2,0)$ LSTs it follows from the M theory definition that the low-energy behaviour (well below the characteristic scale M_s) is given by the low-energy theory on k M5-branes, which is an $\mathcal{N} = (2,0)$ superconformal theory (see, e.g., [7]). For the $(1,1)$ LSTs it follows from the definition using D5-branes in type IIB string theory that the low-energy behaviour is given by an $\mathcal{N} = (1,1)$ $U(k)$ gauge theory whose gauge coupling is $g_{YM}^2 = 1/M_s^2$. For large k , 't Hooft scaling suggests that the perturbative gauge theory (which is free at low energies) breaks down at an energy squared scale of order $1/g_{YM}^2 k = M_s^2/k$.
- (ii) The moduli space metric of theories with 16 supercharges cannot receive any quantum corrections. For the $(1,1)$ LSTs the type IIB constructions show that it is \mathbb{R}^{4k}/S_k , corresponding to the transverse positions of the k identical 5-branes. For the $(2,0)$ LSTs the M theory construction shows that it is $(\mathbb{R}^4 \times S^1)^k/S_k$, where the radius of the S^1 is M_s^2 (recall that the canonical dimension of a scalar field in six dimensions is two). The low-energy theory at generic points in the moduli space involves k tensor multiplets in the $(2,0)$ case and k vector multiplets in the $(1,1)$ case.
- (iii) Type IIA string theory on a circle of radius R is T-dual to type IIB string theory on a circle of radius $1/M_s^2 R$, and an NS 5-brane wrapped on the circle is transformed under this duality to an NS 5-brane wrapped on the dual circle. Thus, T-duality commutes with the limit defining the LSTs, and the $(2,0)$ A_{k-1} LST compactified on a circle of radius R is dual to the type $(1,1)$ A_{k-1} LST compactified on a circle of radius $1/M_s^2 R$. Similarly, the LSTs compactified on T^d have an $O(d, d, \mathbb{Z})$ T-duality symmetry. Note that this shows that T-duality can exist even in non-gravitational theories. The existence of such a T-duality symmetry is the first indication we see that the LSTs are non-local; in particular, after toroidal compactification, they do not have a unique energy-momentum tensor (the same theory can be coupled to different gravitational backgrounds).
- (iv) All LSTs have BPS-saturated strings of tension $T = M_s^2$ (at the origin of moduli space; in some cases there are more BPS-saturated strings away from the origin). These may be viewed as marginally bound states of fundamental strings with the NS 5-branes. At the origin of moduli space, the $(2,0)$ LST has no other BPS states, while in the $(1,1)$ LST the low-energy gluons (and their superpartners) are massless BPS particles. After compactification there are many additional BPS states which we will not discuss here (see, e.g., [3]).

1.2. Motivations

At first sight there is no reason to be interested in “little string theories”, since they are neither directly relevant for studying quantum field theories (being non-local theories) nor for studying quantum gravity (being non-gravitational theories). However, the very fact that these theories are, in some sense, intermediate between local field theories and “standard” string theories means that they may be able to teach us about both. For example, the fact that LSTs have some string theory-like properties, like T-duality and a Hagedorn spectrum (to be discussed below), means that such properties can appear also in non-gravitational theories, and it may be easier to understand them in that context. The existence of non-gravitational Lorentz-invariant theories which are intrinsically non-local is very interesting in itself, and it would be nice to know how to define these theories and what is the proper way to think about them.

Some more concrete motivations for studying these theories are :

- (i) The original construction of LSTs was motivated by the fact that they (or rather their compactification on T^5) arise as the discrete light-cone quantization (DLCQ) of M theory on T^5 [4] with k units of longitudinal momentum.
- (ii) Compactifications of the LSTs lead to many interesting local and non-local field theories. For example, the low-energy limit of the LSTs on T^2 gives four dimensional $\mathcal{N} = 4$ SYM theories, and the compactification of LSTs on other manifolds gives theories related to QCD [8].
- (iii) As we will see in the next section, the LSTs are a rare example of a theory whose discrete light-cone quantization (DLCQ) is simple, and we can use it as a toy model to learn about DLCQ.
- (iv) As we will see in section 3, linear dilaton backgrounds of string theory seem to be holographically dual to LSTs, and thus studying LSTs can teach us about holography in linear dilaton backgrounds (and perhaps give us clues about how to define holography in general backgrounds).

2. DLCQ constructions

The definitions of LSTs given above have not been useful so far for making any explicit computations in these theories, so one would like to have more direct definitions of the LSTs. The first such definition was found using discrete light-cone quantization (DLCQ), which is a quantization of the theory compactified on a light-like circle of radius R with N units of momentum around the compact circle. In the large N limit the momentum around the circle becomes effectively continuous, and it is believed that all six dimensional results may be reproduced. The advantage of DLCQ comes from the fact that negative-momentum modes decouple in the light-cone frame, so the dynamics involves just the positive-momentum modes (for example it can only involve a finite number of particles). To derive a DLCQ theory we need to exactly integrate out the zero modes of fields (carrying no momentum in the compact direction), which is usually

complicated. However, there is a small class of theories where one can derive the exact DLCQ theory, since they have enough supersymmetry to determine the Lagrangian after integrating out the zero modes, and the six dimensional LSTs fall into this class.

The DLCQ of LSTs may be derived either by following Seiberg’s prescription of regarding the DLCQ as the limit of a compactification on a small space-like circle [9], or by starting from M(atric) theory, which is the DLCQ of string theory or M theory, in configurations with 5-branes or singularities, and taking the limit defining the LSTs. All the derivations lead to the same result and they are detailed in the literature (and reviewed in [10]), we will present here only the result.

The DLCQ of the A_{k-1} $(2, 0)$ LST with N units of momentum is [11, 12] the $1 + 1$ dimensional $\mathcal{N} = (4, 4)$ supersymmetric sigma model on the $4Nk$ -dimensional moduli space of N instantons in $SU(k)$ (on \mathbb{R}^4), compactified on a circle of radius $\Sigma = 1/RM_s^2$. Equivalently (by the ADHM construction), it is the conformal theory describing the low-energy limit of the Higgs branch of the $\mathcal{N} = (4, 4)$ $U(N)$ SQCD theory with an adjoint hypermultiplet and k hypermultiplets in the fundamental representation. The moduli space of instantons is singular, but the sigma model on it seems to make sense. The spectrum of chiral operators of the $(2, 0)$ LST may be computed in this description (a similar computation in the $0 + 1$ dimensional sigma model on the moduli space of instantons was performed in [13]).

For the A_{k-1} $(1, 1)$ LSTs it turns out to be complicated to write down the usual DLCQ, but simple to write down a DLCQ with a Wilson line of the low-energy $U(k)$ gauge group around the light-like circle [14, 15, 10]. In the large N limit the effect of the Wilson line is expected to disappear. The DLCQ with the Wilson line is the conformal theory which arises as the low-energy limit of the Coulomb branch of the $1 + 1$ dimensional $\mathcal{N} = (4, 4)$ $U(N)^k$ gauge theory with bifundamental hypermultiplets for adjacent groups (when we arrange the k gauge groups in a circle, as in a quiver diagram). Again, the conformal theory is compactified on a circle of radius $\Sigma = 1/M_s^2 R$. Removing the Wilson line corresponds to taking some of the gauge couplings to infinity before taking the low-energy limit.

The DLCQ gives us an explicit non-perturbative definition of the “little string theories”, enabling us in principle to compute all their states and correlation functions. Unfortunately, in practice such computations are extremely difficult; the conformal theories involved are very complicated even before we take the large N limit, and we have to take this limit to obtain results about the LSTs in six dimensions. Thus, no dynamical computations have been made so far using the DLCQ, but only identifications of some operators and states (see, e.g., [13, 10]). As usual in DLCQ, it is quite complicated to study other points on the moduli space of the LSTs or compactifications, both of which lead to quite different DLCQ theories. For example, the DLCQ of the $(2, 0)$ A_{k-1} theory at a generic point on its moduli space is given by the deformation of the Higgs branch SCFT described above by generic masses for the k hypermultiplets in the fundamental representation (which is a relevant deformation of the Higgs branch SCFT).

3. Holographic constructions

The second useful construction of LSTs is a holographic dual along the lines of the AdS/CFT correspondence; unlike the previous construction this one cannot serve as a definition of the LSTs since we have no non-perturbative definition of the string/M theory backgrounds that are involved, but it seems to be more useful in computing correlation functions in the LSTs (at least at low energies). The AdS/CFT correspondence [16, 18, 17, 19] states that local conformal field theories are often dual to string/M theory compactifications including AdS spaces, and it was generalized to some other classes of local field theories as well (though these are often dual to spaces with regions of high curvature). However, there is no reason for general backgrounds of string/M theory (or any other theory of quantum gravity) to be holographically dual to a local field theory, and it seems likely that the theory which is holographically dual to quantum gravity in Minkowski space is highly non-local. LSTs provide the first example of holography for a non-local non-gravitational theory, and it turns out that generally LST-like theories are holographically dual to backgrounds which asymptote to string theory in a linear dilaton background [20].

As in the AdS/CFT correspondence, the space which is holographically dual to LSTs may be derived by starting from the string theory background corresponding to NS 5-branes and taking the limit of $g_s \rightarrow 0$ (where g_s is the asymptotic string coupling) which defines the LSTs [20, 21, 22]. For $(2, 0)$ LSTs the correct procedure is actually to start from the background corresponding to M5-branes with a transverse circle, as described above, since a background of NS 5-branes in string theory does not correspond to a configuration which is localized in the S^1 coordinates of the moduli space. Starting with such a background and taking the appropriate limit leads to the space

$$l_p^{-2} ds^2 = H^{-1/3} dx_6^2 + H^{2/3} (dx_{11}^2 + dU^2 + U^2 d\Omega_3^2), \quad (3.1)$$

where dx_6^2 is the metric on \mathbb{R}^6 , $d\Omega_3^2$ is the metric on S^3 , x_{11} is compactified on a circle with radius M_s^2 ,

$$H = \sum_{j=-\infty}^{\infty} \frac{\pi k}{(U^2 + (x_{11} - 2\pi j M_s^2)^2)^{3/2}}, \quad (3.2)$$

and there are also k units of 4-form flux on $S^3 \times S^1$. M theory compactified on this space is holographically dual to the $A_{k-1} (2, 0)$ LST.

The space (3.1) is quite complicated, but it simplifies in the asymptotic regions of space. For small U and x_{11} the space (3.1) becomes just $AdS_7 \times S^4$, which is holographically dual to the $(2, 0)$ SCFTs, as expected since these SCFTs are the low-energy limit of the $(2, 0)$ LSTs. For large U the physical radius of the x_{11} circle becomes very small and it is more appropriate to view the background as a type IIA string theory compactification. Defining a new variable ϕ by $l_s^2 U = \sqrt{k} e^{\phi/\sqrt{k} l_s}$, the string metric is simply

$$ds_{string}^2 = dx_6^2 + d\phi^2 + k l_s^2 d\Omega_3^2, \quad (3.3)$$

with a linear dilaton in the ϕ direction,

$$g_s(\phi) = e^{-\phi/\sqrt{k}l_s}, \quad (3.4)$$

and k units of 3-form flux on the S^3 . Note that for large k the curvatures are small everywhere, so supergravity is a good approximation at low energies. Even though the behaviour of this space near the boundary is very different from that of AdS space, it seems that the same principles apply also in this case; for example, as in AdS/CFT, correlation functions in the LSTs are identified with the response of string/M theory on the background (3.1) to turning on boundary values for non-normalizable modes of the fields. In the space (3.1) some of these modes are just the incoming and outgoing waves in the ϕ direction, so that some correlation functions of the LST correspond to the S-matrix for scattering in the ϕ direction.

Similarly, the type (1,1) LSTs are holographically dual to the near-horizon limit of type IIB NS 5-branes; unfortunately this background becomes singular for small U (there is no analog of the M theory region in (3.1)) so this description is less useful for computations of correlation functions, but it can still be used for identifying the operators and some states of these LSTs as described below.

The holographic description of the LSTs is useful for :

- (i) As in the AdS/CFT correspondence, some correlation functions may be computed using supergravity. In this case the supergravity approximation is valid for the (2,0) LSTs at large k and at energies well below the string scale M_s , where stringy corrections in the background (3.3) become important. The computation of the 2-point function of the energy-momentum tensor in the LSTs was described in [23], and it is possible to compute also other correlation functions (though the actual computations are quite complicated). In particular, supergravity gives an exact description of the analog of the 't Hooft limit for the (2,0) LSTs, in which M_s and k are taken to infinity with the scale M_s^2/k (which in the (1,1) case is the inverse 't Hooft gauge coupling) kept constant.
- (ii) A big difference between the background (3.1) and backgrounds which are dual to other field theories is that near the boundary of (3.1) the string coupling goes to zero and the curvatures are small, so one can compute the spectrum of fields exactly (and not just for large k as in other cases). The full spectrum of chiral fields in the LSTs was computed in this way in [20], and turned out to be exactly the same for all values of k (in the A_{k-1} case) as the spectrum of chiral fields in the field theories which arise as the low-energy limit of the LSTs.
- (iii) The holographic description can be used to reliably compute some of the states in the LSTs, which are states propagating in the weakly coupled region of (3.1). For example, one has the states in the supergraviton multiplet propagating with some momentum in the ϕ direction, and these look like a continuum of states from the six dimensional point of view (the mass shell condition relates the 6-dimensional momentum to the ϕ -momentum but it does not determine its magnitude). It turns out that for the states in the supergravity multiplet this continuum of states starts

at a scale of order $M^2 \simeq M_s^2/k$, while other string states in the weak coupling region also give rise to continuous spectra in the LST, starting at higher values of M^2 . The implications of the existence of this continuum are not clear. The same continuum of states can also be seen in the DLCQ formalism [10], where it is related to the continuum of states found in sigma models on orbifold singularities with zero theta angle.

- (iv) The holographic description can be used to analyze the behaviour at finite temperature and energy density, as discussed in the next section.

4. Behaviour at finite energy density

We have already seen some indications that LSTs are not local field theories, but the strongest indication for this seems to come from analyzing the equation of state of the LSTs at finite temperature or energy density. For local field theories the high-energy behaviour of the density of states is always an exponential of a power of the energy density which is less than one, while for the LSTs we will see that it is exponential in the energy density.

The most reliable way to compute the equation of state of the LSTs is to use their holographic description. Holographic dualities relate finite temperature states of non-gravitational theories to black hole configurations, with the Hawking temperature of the black hole equated with the field theory temperature, and the field theory energy equated with the mass of the black hole. The entropy of such black holes can be computed by the usual Bekenstein-Hawking formula which relates it to their area. In the case of LSTs the appropriate black holes (at large enough energy densities) are the near-horizon limits of near-extremal NS5-branes and this computation was first done in this context in [24]. The result is

$$E = T_H S; \quad T_H \simeq M_s / \sqrt{6k}, \quad (4.1)$$

and it is reliable when the curvatures in the black hole background are small (requiring $k \gg 1$) and when the string coupling at the black hole horizon is small (recall that the string coupling becomes weaker and weaker as one goes towards the boundary), requiring that the energy density $\mu \equiv E/V$ satisfies $\mu \gg kM_s^6$. Thus, we find that for large k and large energy densities the equation of state is Hagedorn-like, corresponding to a limiting temperature in the theory of $T_H \simeq M_s / \sqrt{6k}$; above some energy density which is smaller than kM_s^6 the specific heat diverges and increasing the energy density no longer increases the temperature, which remains $T = T_H$.

This behaviour is similar to the behaviour of single-string states in free string theory, where $T_H \simeq M_s$; however, in string theory it is believed (at least in some cases) that at the temperature T_H there is a phase transition to a different phase where a state becoming massless at $T = T_H$ condenses [25], while in the LSTs there is no evidence for any states becoming massless at $T = T_H$, so it seems that T_H is really a limiting temperature (and the canonical ensemble cannot be defined beyond this temperature).

It would be nice to reproduce the behaviour (4.1) also in the DLCQ description of the LSTs, but this seems rather difficult. The behaviour (4.1) is, in fact, reproduced by a naive counting of the density of states in the DLCQ description [26], but the relevant states are really only those whose energy scales as $1/N$ in the large N limit, and identifying all these states is much more difficult. Some states whose energy scales as $1/N$ were identified in [10], and they correspond to a density of states in space-time which is exponential but with a smaller exponent corresponding to $\tilde{T}_H \simeq M_s/\sqrt{12}$; it would be interesting to identify in the DLCQ the other states which give rise to (4.1).

It is not clear how to interpret the exponential behaviour of the equation of state in the LSTs. The fact that the density of states depends only on the energy and not on the volume suggests that generic high-energy states are single-object configurations, perhaps similar to long, space-filling strings. It is not yet known at which energy density the behaviour (4.1) begins, and what is the equation of state in the regime $\mu < M_s^6$ (which is relevant, in particular, for the ’t Hooft limit of the LSTs described above).

General arguments (with some assumptions about the behaviour of generic operators in the LSTs) suggest that the behaviour (4.1) implies the non-existence of correlation functions of operators at separations smaller than T_H^{-1} . Presumably, there are no local operators in these theories, but only operators “smeared” on distance scales of at least T_H^{-1} . Computations in the holographic formulation give correlation functions of operators (which are naively local) in momentum space, but it seems that the Fourier transform of the results to position space does not exist for small separations [27], consistent with this non-locality. Note that the scale of non-locality suggested by these arguments is different (by a factor of order \sqrt{k}) from the scale of non-locality suggested by the T-duality symmetry.

5. Future directions

Much progress has been made in the last few years in understanding “little string theories”, but many open questions remain. Some interesting open questions are :

- (i) The behaviour of LSTs at high energies is clearly not governed by a field theoretical fixed point, and it is interesting to understand what the high-energy behaviour is. There are some indications that the theory becomes weakly coupled at high energies, such as the fact that in the holographic description the string coupling vanishes near the boundary, and that the string coupling is small everywhere in the holographic dual of the theory at high energy densities. It would be interesting to understand if there is some sense in which the theory becomes weakly coupled at high energies (perhaps as in asymptotically free field theories) and if there is any simple description of the high-energy limit.
- (ii) The behaviour at intermediate energy scales, such as those governed by the ’t Hooft limit (where we take large k and look at energies of order $E \sim M_s/\sqrt{k}$), is still not clear. In principle correlation functions at these scales may be computed from supergravity, but so far it is not known if their behaviour (and the behaviour of the

equation of state in this regime) is more like a field theory (and, if so, which field theory ?) or more like a string theory.

- (iii) We have not had a chance to discuss here LSTs in lower dimensions or with less supersymmetry. It would be interesting to try to classify these theories and to see if they are similar to the six dimensional LSTs described above or not. It is also interesting to analyze compactifications of the LSTs, which give rise to many interesting theories. Unlike local field theories, the behaviour of non-local field theories upon compactification is not determined by the behaviour of the uncompactified theory, so it should be studied independently and may reveal new results. In particular, the large k behaviour of the T^5 compactification of the LSTs is related to M(atric) theory on T^5 , and it would be interesting to understand it further and to identify the states whose energy scales as $1/k$ in the large k limit.

It has recently been suggested [28, 29, 30] that the holographic description of LSTs becomes weakly coupled if one looks at particular configurations which are far out on the LST moduli space. For example, configurations in the $(1, 1)$ theory where the low-energy $U(k)$ gauge theory is broken at a scale M_W were argued to have a perturbation expansion in M_s/M_W , which can be used to reliably compute correlation functions in these theories (far out on the moduli space) at energies well below M_W (but potentially above the scales M_s and M_s/\sqrt{k}). If perturbation theory is indeed reliable in these configurations it would be very interesting to use them for various computations, and to see what they can teach us about the LSTs.

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